

Resonances in weak nonleptonic Ω^- decay

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Abstract

We examine the importance of $J^P = \frac{1}{2}^+, \frac{1}{2}^-$ resonances for weak nonleptonic Ω^- decays within the framework of chiral perturbation theory. The spin-1/2 resonances are included into an effective theory and tree contributions to the Ω^- decays are calculated. We find significant contributions to the decay amplitudes and satisfactory agreement with experiment. This confirms and extends previous results wherein such spin-1/2 resonances were included in nonleptonic and radiative-nonleptonic hyperon decays.

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1 Introduction

Among the $J^P = \frac{3}{2}^+$ decuplet states only the Ω^- decays weakly into nonleptonic channels. The allowed two-body decay modes to the lowest-lying baryon octet are $\Omega^- \rightarrow \Lambda K^-$, $\Omega^- \rightarrow \Xi^0 \pi^-$ and $\Omega^- \rightarrow \Xi^- \pi^0$. The pertinent decay amplitudes have been calculated within the framework of heavy baryon chiral perturbation theory at tree level [1] and also at one-loop order including the lowest nonanalytic corrections.[2] Within these investigations the authors come to the conclusion that the decay amplitudes are well described by values for the weak parameters of the Lagrangian determined from a fit to the s-wave amplitudes of the nonleptonic decays of the octet baryons. But the same values inserted into the expressions for the p-wave amplitudes of the nonleptonic hyperon decays do not allow a good description of the experimental data.[1] Therefore, it is questionable if the chosen values for the weak parameters are appropriate.

An intriguing approach for nonleptonic hyperon decays was examined by Le Yaouanc et al., who assert that a reasonable fit for both s- and p-waves can be provided by appending pole contributions from $SU(6)$ $(70, 1^-)$ states to the s-wave amplitudes.[3] Their calculations were performed in a simple constituent quark model and appeared to be able to provide a resolution of the s- and p-wave dilemma. In a previous work [4] we have studied this approach within a chiral framework and have shown that it is indeed possible to find a simultaneous fit to both s- and p-wave hyperon decay amplitudes if contributions from the lowest lying $1/2^-$ and $1/2^+$ baryon octet resonant states are included in the formalism.

Other weak processes involving hyperons are, e.g., the radiative-nonleptonic hyperon decays. The primary problem of radiative hyperon decay has been to understand the large negative value found experimentally for the asymmetry parameter in polarized $\Sigma^+ \rightarrow p \gamma$ decay.[5] The difficulty here is associated with the restrictions posed by Hara's theorem, which requires the vanishing of this asymmetry in the $SU(3)$ limit.[6] Recent work involving the calculation of chiral loops has also not lead to a resolution, although slightly larger asymmetries can be accommodated.[7] In [8] we examined the contribution of $1/2^-$ and $1/2^+$ baryon intermediate states to radiative hyperon decay within the framework of chiral perturbation theory. We obtained reasonable predictions for the decay amplitudes and a significant negative value for the $\Sigma^+ \rightarrow p \gamma$ asymmetry as a very natural result of this picture, even though Hara's theorem is satisfied. Thus, the inclusion of spin-1/2 resonances provides a reasonable explanation of the importance of higher order counterterms and gives a satisfactory picture of both radiative and nonradiative nonleptonic hyperon decays.

In the present paper we extend the discussion to weak nonleptonic Ω^- transitions and investigate the validity of this approach for these decays. In the next Section we introduce the effective weak and strong Lagrangian including resonant states and evaluate the pole diagram contributions of both ground state baryons and resonant states to nonleptonic Ω^- decay. Numerical results are presented in Sec. 3, and in Sec. 4 we conclude with a short summary. In the Appendix, we determine the strong couplings of the decuplet to the spin-1/2 resonances by a fit to the strong decays of these spin-1/2 resonances.

2 Weak nonleptonic Ω^- decay

There are three two-body decay modes of the Ω^- to the ground state baryon octet: $\Omega^- \rightarrow \Lambda K^-$, $\Omega^- \rightarrow \Xi^0 \pi^-$ and $\Omega^- \rightarrow \Xi^- \pi^0$. Phenomenologically, the matrix elements of these decays can each be expressed in terms of a parity-conserving p-wave amplitude $A_{ij}^{(P)}$ and a parity-

violating d-wave amplitude $A_{ij}^{(D)}$

$$\mathcal{A}(\Omega^- \rightarrow B_i \phi_j) = \bar{u}_{B_i} \left\{ A_{ij}^{(P)} q_\mu + A_{ij}^{(D)} \gamma_5 q_\mu \right\} u_{\Omega^-}^\mu, \quad (1)$$

where B_i and ϕ_j are the ground state baryon and Goldstone boson, respectively, and q_μ is the four-momentum of the outgoing kaon or pion. The underlying strangeness-changing Hamiltonian transforms under $SU(3)_L \times SU(3)_R$ as $(8_L, 1_R) \oplus (27_L, 1_R)$ and, experimentally, the octet piece dominates over the 27-plet by a factor of twenty or so in nonleptonic hyperon decay. In the case of Ω^- decay the corresponding octet dominance contribution implies the relation

$$\mathcal{A}(\Omega^- \rightarrow \Xi^0 \pi^-) - \sqrt{2} \mathcal{A}(\Omega^- \rightarrow \Xi^- \pi^0) = 0 \quad (2)$$

which holds both for p- and d-waves. The experimentally measured ratio of the decay widths — $\Gamma(\Omega^- \rightarrow \Xi^0 \pi^-)/\Gamma(\Omega^- \rightarrow \Xi^- \pi^0) = 2.65 \pm 0.20$ — is in disagreement with the isospin prediction of 2. This indicates that the $\Delta I = 3/2$ amplitude in $\Gamma(\Omega^- \rightarrow \Xi \pi)$ may be significantly larger than in nonleptonic hyperon decay [9] and could signal a violation of the $\Delta I = 1/2$ rule for the Ω^- decay.[10] We do not study this issue here and have, therefore, neglected the 27-plet contribution. We prefer not to work in the isospin limit, however, and do not use this isospin relation to eliminate one of the decay amplitudes but rather attempt to find a simultaneous fit to the three decay modes of the Ω^- .

The purpose of this work is to study the role of spin-1/2 resonances in Ω^- decays. To this end, it is sufficient to work at tree level. This is also necessary because the calculations [4, 8] which provide values for the weak parameters involved in these decays have been performed at tree level, so that it would be inconsistent to use these values for the weak parameters in a loop calculation for the Ω^- decays. In the rest frame of the Ω^- the amplitudes are related to the decay width via

$$\Gamma = \frac{1}{12\pi} \frac{|\mathbf{q}|}{M_\Omega} (E_\phi^2 - m_\phi^2) \left[|A^{(P)}|^2 (E_B + M_B) + |A^{(D)}|^2 (E_B - M_B) \right] \quad (3)$$

with \mathbf{q} being the three-momentum of the outgoing pseudoscalar. In previous work the calculation of the d-wave amplitudes $A^{(D)}$ has generally been neglected, since its contribution to the decay width is suppressed by kinematical factors. However, in order to evaluate the decay parameters on the other hand, one *must* include the d-waves. Thus, e.g., the asymmetry parameter is

$$\alpha = \frac{2\text{Re}(A^{(P)*} \bar{A}^{(D)})}{|A^{(P)}|^2 + |\bar{A}^{(D)}|^2} \quad (4)$$

with $\bar{A}^{(D)} = [(E_B - M_B)/(E_B + M_B)]^{1/2} A^{(D)}$. [5]

2.1 The effective Lagrangian

The effective Lagrangian can be decomposed into a strong and weak component

$$\mathcal{L}_{\text{eff}} = \mathcal{L}^{(S)} + \mathcal{L}^{(W)}. \quad (5)$$

As mentioned above, since we will be using results from our previous work, which considered the role of these resonances in nonleptonic hyperon decay at tree level, we restrict ourselves to

tree level for the present study. We first consider the Lagrangian in the absence of spin-1/2 resonances. The strong part consists of the free kinetic Lagrangians of the decuplet, the ground state baryon octet and the Goldstone bosons together with an interaction term

$$\mathcal{L}^{(S)} = \mathcal{L}_{kin} + \mathcal{L}_{\Delta B\phi}. \quad (6)$$

The kinetic component reads

$$\begin{aligned} \mathcal{L}_{kin} = & i \text{tr}(\bar{B}\gamma_\mu[D^\mu, B]) - \overset{\circ}{M} \text{tr}(\bar{B}B) + \frac{F_\pi^2}{4} \text{tr}(u_\mu u^\mu) + \frac{F_\pi^2}{4} \text{tr}(\chi_+) \\ & + \bar{\Delta}^\alpha \left[(-i\not{\partial} + M_\Delta)g_{\alpha\beta} + i(\gamma_\alpha\partial_\beta + \gamma_\beta\partial_\alpha) - i\gamma_\alpha\not{\partial}\gamma_\beta - M_\Delta\gamma_\alpha\gamma_\beta \right] \Delta^\beta \end{aligned} \quad (7)$$

with $\overset{\circ}{M}$ being the octet baryon mass in the chiral limit. The pseudoscalar Goldstone fields ($\phi = \pi, K, \eta$) are collected in the 3×3 unimodular, unitary matrix $U(x)$,

$$U(\phi) = u^2(\phi) = \exp\{2i\phi/F_\pi\} \quad , \quad u_\mu = iu^\dagger \nabla_\mu U u^\dagger \quad (8)$$

where $F_\pi \simeq 92.4$ MeV is the pion decay constant,

$$\phi = \frac{1}{\sqrt{2}} \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{6}}\eta & K^0 \\ K^- & \bar{K}^0 & -\frac{2}{\sqrt{6}}\eta \end{pmatrix} \quad (9)$$

represents the contraction of the pseudoscalar fields with the Gell-Mann matrices and B is the standard $SU(3)$ matrix representation of the low-lying spin-1/2 baryons (N, Λ, Σ, Ξ). Here $\chi_+ = u^\dagger \chi u^\dagger + u \chi^\dagger u$ is proportional to the quark mass matrix $\mathcal{M} = \text{diag}(m_u, m_d, m_s)$ via $\chi = 2B\mathcal{M}$. Also, $B = -\langle 0|\bar{q}q|0\rangle/F_\pi^2$ is the order parameter of the spontaneous symmetry violation, and we assume $B \gg F_\pi$. The propagator of the spin-3/2 fields (denoted generically by Δ) is given by

$$G_{\beta\delta}(p) = -i \frac{\not{p} + M_\Delta}{p^2 - M_\Delta^2} \left(g_{\beta\delta} - \frac{1}{3} \gamma_\beta \gamma_\delta - \frac{2p_\beta p_\delta}{3M_\Delta^2} + \frac{p_\beta \gamma_\delta - p_\delta \gamma_\beta}{3M_\Delta} \right) , \quad (10)$$

with $M_\Delta = 1.38$ GeV being the average decuplet mass.

The interaction Lagrangian between the spin-3/2 fields, the baryon octet and the Goldstone bosons reads

$$\mathcal{L}_{\Delta B\phi} = \frac{\mathcal{C}}{2} \left\{ \bar{\Delta}^{\mu,abc} \Theta_{\mu\nu}(Z) (u^\nu)_a^i B_b^j \epsilon_{cij} - \bar{B}_i^b (u^\nu)_j^a \Theta_{\nu\mu}(Z) \Delta_{abc}^\mu \epsilon^{cij} \right\} , \quad (11)$$

where a, b, \dots, j are $SU(3)_f$ indices and the coupling constant $1.2 < \mathcal{C} < 1.8$ has been determined from the decays $\Delta \rightarrow B\pi$. [11] In our application, we use the mean value $\mathcal{C} = 1.5$. The Dirac matrix operator $\Theta_{\mu\nu}(Z)$ is given in general by

$$\Theta_{\mu\nu}(Z) = g_{\mu\nu} - \left(Z + \frac{1}{2} \right) \gamma_\mu \gamma_\nu . \quad (12)$$

However, the off-shell parameter Z does not contribute at tree level because of the subsidiary condition $\gamma_\mu u_{\Omega-}^\mu = 0$.

The weak Lagrangian can be written as

$$\mathcal{L}^{(W)} = d \operatorname{tr}(\bar{B}\{h_+, B\}) + f \operatorname{tr}(\bar{B}[h_+, B]) + \frac{F_\pi^2}{4} h_\pi \operatorname{tr}(h_+ u_\mu u^\mu) + h_c \bar{\Delta}^{\mu, abc} (h_+)_a^d \Delta_{\mu, dbc}, \quad (13)$$

where we have defined

$$h_+ = u^\dagger h u + u^\dagger h^\dagger u, \quad (14)$$

with $h_b^a = \delta_2^a \delta_b^3$ being the weak transition matrix. (Note that h_+ transforms as a matter field.) In Eq.(13) the weak coupling h_π is known from weak nonleptonic kaon decays — $h_\pi = 3.2 \times 10^{-7}$ — while in our previous work the LECs d and f have been determined from nonleptonic hyperon decay: $d = 0.44 \times 10^{-7}$ GeV, $f = -0.50 \times 10^{-7}$ GeV. In [4] we included lowest lying spin-1/2⁺ and 1/2[−] resonances in the theory and performed a tree level calculation. Integrating out the heavy degrees of freedom provides then a plausible estimate of the weak counterterms and a satisfactory fit for both s- and p-waves was achieved. Since we herein apply this scheme for the related weak Ω^- decays we will use the values for d and f from [4]. The parameter h_c does not appear in the tree-level calculation of the nonleptonic hyperon decays in [4, 8]. We therefore consider it as a free parameter and determine its value from a fit to the weak Ω^- decays.

We now proceed to append the contribution from spin-1/2 resonances. In [3, 12] it was argued that in a simple constituent quark model inclusion of the lowest lying spin 1/2[−] octet from the (70,1[−]) multiplet leads to significant improvements in both radiative and nonleptonic hyperon decays, and we confirmed in two recent calculations [4, 8] that there indeed exist significant contributions from such resonances to the nonleptonic hyperon decays in the framework of chiral perturbation theory. We begin therefore with the inclusion of the octet of spin-parity 1/2[−] states, which include the well-established states $N(1535)$ and $\Lambda(1405)$. As for the rest of the predicted 1/2[−] states there are a number of not so well-established states in the same mass range — cf. [3] and references therein. In order to include such resonances one begins by writing down the most general Lagrangian at lowest order which exhibits the same symmetries as the underlying theory, i.e. Lorentz invariance and chiral symmetry. For the strong part we require invariance under C and P transformations separately, while the weak piece is invariant under CPS transformations, where the transformation S interchanges down and strange quarks in the Lagrangian. We will work in the CP -conserving limit so that all LECs are real, and denote the 1/2[−] octet by R . Another important multiplet of excited states is the octet of Roper-like spin-1/2⁺ fields. While it was argued in [11] that these play no significant role, a more recent study seems to indicate that one cannot neglect contributions from such states to, e.g., decuplet magnetic moments.[13] Also in [4, 8] we found that inclusion of these states was essential both in nonleptonic and radiative hyperon decay to achieve experimental agreement. It is thus important to include also the contribution of these baryon resonances here. The Roper octet, which we denote by B^* , consists of the $N^*(1440)$, the $\Sigma^*(1660)$, the $\Lambda^*(1600)$ and the $\Xi^*(1620?)$.

The resonance kinetic term is straightforward

$$\mathcal{L}_{kin} = i \operatorname{tr}(\bar{R} \gamma_\mu [D^\mu, R]) - M_R \operatorname{tr}(\bar{R} R) + i \operatorname{tr}(\bar{B}^* \gamma_\mu [D^\mu, B^*]) - M_{B^*} \operatorname{tr}(\bar{B}^* B^*) \quad (15)$$

with M_R and M_{B^*} being the masses of the resonance octets in the chiral limit. The strong interaction Lagrangian relevant to the processes considered here reads

$$\mathcal{L}_{\Delta\phi B^*/R} = \frac{\mathcal{C}_{B^*}}{2} \left\{ \bar{\Delta}^{\mu, abc} \Theta_{\mu\nu}(Z) (u^\nu)_a^i B_b^{*j} \epsilon_{cij} - \bar{B}_i^{*b} (u^\nu)_j^a \Theta_{\nu\mu}(Z) \Delta_{abc}^\mu \epsilon^{cij} \right\}$$

$$+ i \frac{s_c}{2} \left\{ \bar{\Delta}^{\mu, abc} \Theta_{\mu\nu}(Z) (u^\nu)_a^i \gamma_5 R_b^j \epsilon_{cij} - \bar{R}_i^b (u^\nu)_j^a \Theta_{\nu\mu}(Z) \gamma_5 \Delta_{abc}^\mu \epsilon^{cij} \right\} \quad (16)$$

and the couplings C_{B^*} and s_c can be determined from a fit to the strong decays of the spin-1/2 resonances to $\Delta(1232)$ — cf. App. A. The corresponding weak Lagrangian is

$$\begin{aligned} \mathcal{L}^W = & d^* \left[\text{tr}(\bar{B}^* \{h_+, B\}) + \text{tr}(\bar{B} \{h_+, B^*\}) \right] + f^* \left[\text{tr}(\bar{B}^* [h_+, B]) + \text{tr}(\bar{B} [h_+, B^*]) \right] \\ & + iw_d \left[\text{tr}(\bar{R} \{h_+, B\}) - \text{tr}(\bar{B} \{h_+, R\}) \right] + iw_f \left[\text{tr}(\bar{R} [h_+, B]) - \text{tr}(\bar{B} [h_+, R]) \right] \end{aligned} \quad (17)$$

with four couplings d^*, f^* and w_d, w_f which have been determined from a fit to the nonleptonic hyperon decays in [4].

There is thus only a single unknown parameter in this approach — h_c — once the weak couplings are fixed from nonleptonic hyperon decays. The inclusion of the spin-1/2 resonant states does not lead to any new unknown parameters, so that study of the weak nonleptonic Ω^- decays provides a nontrivial check on the role of spin-1/2 resonances in weak baryon decays — extending this approach to the case of weak decuplet decays.

2.2 Pole contributions

In this section we consider the tree diagrams which contribute to weak Ω^- decay. The graphs involving only the decuplet and the ground state baryon octet are depicted in Fig. 1. Graphs 1a and 1b contribute to the decay $\Omega^- \rightarrow \Lambda K^-$, while graphs 1b and 1c deliver contributions to the pionic decays $\Omega^- \rightarrow \Xi \pi$. (Note that diagram 1c with the weak decay of the Goldstone boson has been neglected in [1] and [2].) The diagrams in Fig. 1 contribute only to the parity-conserving p-wave amplitudes $A^{(P)}$, yielding

$$\begin{aligned} A_{\Lambda K}^{(P)} &= \frac{C}{2\sqrt{3}F_\pi} \left(\frac{d - 3f}{M_\Lambda - M_{\Xi^0}} + \frac{h_c}{M_\Omega - M_{\Xi^*}} \right) \\ A_{\Xi^0 \pi^-}^{(P)} &= \frac{C}{\sqrt{2}F_\pi} \left(\frac{h_c}{3(M_\Omega - M_{\Xi^*})} + \frac{h_\pi m_{\pi^-}^2}{2(m_{K^-}^2 - m_{\pi^-}^2)} \right) \\ A_{\Xi^- \pi^0}^{(P)} &= \frac{C}{2F_\pi} \left(\frac{h_c}{3(M_\Omega - M_{\Xi^*})} + \frac{h_\pi m_{\pi^0}^2}{2(m_{K^0}^2 - m_{\pi^0}^2)} \right). \end{aligned} \quad (18)$$

Here, we have used the physical meson masses and have replaced the baryon masses in the chiral limit by their physical values, which is consistent to the order we are working. There are no contributions from these diagrams to the parity-violating d-wave amplitudes $A^{(D)}$, so that at this order there is in each case a vanishing asymmetry parameter α .

We now include the spin-1/2 resonances, which contribute only to the decay $\Omega^- \rightarrow \Lambda K^-$ through the diagram depicted in Fig. 2. One obtains the result

$$\begin{aligned} A_{\Lambda K}^{(P)} &= \frac{C_{B^*}}{2\sqrt{3}F_\pi} \frac{d^* - 3f^*}{M_\Lambda - M_{B^*}} \\ A_{\Lambda K}^{(D)} &= \frac{s_c}{2\sqrt{3}F_\pi} \frac{w_d - 3w_f}{M_\Lambda - M_R} \end{aligned} \quad (19)$$

and there are no further contributions from the spin-1/2 resonant states.

3 Numerical results

In this section we present the results obtained by a fit to the three weak Ω^- decays. (Note that when applying Eq.(3) in order to perform the fit, we use the physical values for the masses of the outgoing particles. Therefore, we obtain a ratio $\Gamma(\Omega^- \rightarrow \Xi^0 \pi^-)/\Gamma(\Omega^- \rightarrow \Xi^- \pi^0)$ which is slightly different from the isospin prediction of 2. However, the 30% effect found experimentally is outside our approach.) Our fit including the spin-1/2 resonances leads to $h_c = 0.39 \times 10^{-7}$ GeV and in units of 10^{-15} GeV

$$\begin{aligned}\Gamma(\Omega^- \rightarrow \Lambda K^-) &= 5.61 & (5.42 \pm 0.06) \\ \Gamma(\Omega^- \rightarrow \Xi^0 \pi^-) &= 1.62 & (1.89 \pm 0.06) \\ \Gamma(\Omega^- \rightarrow \Xi^- \pi^0) &= 0.76 & (0.69 \pm 0.03),\end{aligned}\tag{20}$$

where we have given the experimental numbers in brackets. (Note that the ratio $\Gamma(\Omega^- \rightarrow \Xi^0 \pi^-)/\Gamma(\Omega^- \rightarrow \Xi^- \pi^0)$ is approximately 2.1 in our approach.) The inclusion of the spin-1/2 resonances predicts a finite decay parameter α for the decay $\Omega^- \rightarrow \Lambda K^-$

$$\alpha_{\Lambda K} = -0.015\tag{21}$$

to be compared with an experimental value of -0.026 ± 0.026 , while we obtain vanishing asymmetry parameters for the two pionic decay modes which are also consistent with the experimental values $\alpha_{\Xi^0 \pi^-} = 0.09 \pm 0.14$ and $\alpha_{\Xi^- \pi^0} = 0.05 \pm 0.21$.³

Our results are only indicative. A full discussion would have to include both the effects of chiral loops as well as contributions from higher resonances. In this regard, we do not anticipate our results to be able to precisely reproduce the experimental values for the decay widths and the asymmetry parameters. Although it would be desirable to estimate the error of our tree level result from higher order chiral loop corrections by e.g. calculating some typical loop contributions, it is then inconsistent to use results from our previous work which considered the role of resonances in nonleptonic hyperon decay at tree level [4, 8]. Therein we derived estimates on the resonance parameters at tree level which we also use in the present investigation. The inclusion of loops might change the numerical values of these parameters somewhat and, therefore, alter the contribution from resonances both in nonleptonic hyperon decay and Ω^- decay. In order to give a reliable estimate on chiral corrections in Ω^- decay, one has to consider—in addition to chiral loops for this process—also the change in the resonance contribution due to the inclusion of loops in nonleptonic hyperon decay. This is clearly beyond the scope of the present investigation. Our purpose herein is to examine whether the inclusion of spin-1/2 resonances, which worked well in both ordinary and radiative nonleptonic hyperon decay, can be extended successfully to the case of weak decuplet decays. Our results suggest that the spin-1/2 resonances also play an essential role in weak nonleptonic decuplet decay.

Before leaving this section, it is interesting to note that while our results are purely phenomenological, it is possible to compare the fit value of h_c , which leads in our model to $\langle \Xi^{*-} | H_W | \Omega^- \rangle =$

³The results from a fit without the inclusion of the spin-1/2 resonances read in units of 10^{-15} GeV $\Gamma(\Omega^- \rightarrow \Lambda K) = 4.64, \Gamma(\Omega^- \rightarrow \Xi^0 \pi^-) = 0.69, \Gamma(\Omega^- \rightarrow \Xi^- \pi^0) = 0.32$ with $h_c = 0.24 \times 10^{-7}$ GeV and vanishing asymmetry parameters in each case. Interestingly, we do not have a satisfactory fit to experiment. However, it is not really fair to compare these results to the fit including the spin-1/2 resonances, since we employ values for d and f which have been determined in [4] from a fit to the nonleptonic decays after the inclusion of such states.

0.22×10^{-7} GeV, with what is expected based upon quark model considerations. A constituent quark model approach to the weak $\Xi^{*-} - \Omega^-$ matrix element would include two structures — one from the direct four-quark weak Hamiltonian and a second from the related penguin term. However, the direct contribution to the $\Xi^{*-} - \Omega^-$ transition vanishes because of lack of multiple s, \bar{s} quarks in H_W , so that a non-zero value for this matrix element can only arise from the penguin term. The size of this contribution is somewhat larger than what one might expect based on model calculations with perturbatively calculated Wilson coefficients evolved to $\mu \simeq 1$ GeV, but is consistent with the expected size of such effects when scaled into the low energy region $\mu \simeq 0.2$ GeV.[14]

4 Summary

In this work we have examined the role of spin-1/2 intermediate state resonances in weak nonleptonic Ω^- decay. This study also serves as a consistency check for the results from nonleptonic hyperon decay. In two recent papers we were able to show that these resonances play an essential role for both ordinary and radiative nonleptonic hyperon decay.[4, 8] A much improved agreement with the experimental values for s- and p-waves in ordinary hyperon decay is brought about even at tree level and a significant negative value for the asymmetry parameter in the radiative decay $\Sigma^+ \rightarrow p\gamma$ is obtained as a very natural result within this picture.

In the present work we have shown that this approach can be extended to the case of weak decuplet decays. The pertinent Lagrangian without spin-1/2 resonances has a single unknown weak parameter once the remaining weak parameters are fixed by the nonleptonic hyperon decays. Inclusion of the spin-1/2 resonances does not lead to any additional unknown parameters, since the two strong couplings involving the decuplet can be fixed from the strong decays of the spin-1/2 resonances to the decuplet. The spin-1/2 resonances contribute only to the decay $\Omega^- \rightarrow \Lambda K^-$. We determine the unknown parameter by a least-squares fit to the three two body decay modes of the Ω^- . One obtains satisfactory agreement with experimental data and a nonvanishing asymmetry parameter $\alpha_{\Lambda K} = -0.014$ which lies within the experimental range — $\alpha_{\Lambda K} = -0.026 \pm 0.026$. Since we do not work in the isospin limit we obtain a ratio $\Gamma(\Omega^- \rightarrow \Xi^0 \pi^-)/\Gamma(\Omega^- \rightarrow \Xi^- \pi^0) = 2.1$ which is different from the isospin prediction of 2. This ratio is measured to be approximately 2.65.[5] This indicates that the $\Delta I = 3/2$ amplitude from the current measurement of the rates for $\Gamma(\Omega^- \rightarrow \Xi \pi)$ appears to be significantly larger than in nonleptonic hyperon decay [9] and could signal a violation of the $\Delta I = 1/2$ rule for the Ω^- decay.[10] We do not test that here and have, therefore, neglected the 27-plet contribution.

Our study suggests that the approach of including spin-1/2 resonances in nonleptonic hyperon decay can be extended to weak decuplet decays giving a satisfactory picture of weak Ω^- decay. In order to make a more definite statement one should, of course, go to higher orders and include meson loops as well as the contributions from additional resonances. However, this is beyond the scope of the present investigation.

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A Determination of the spin-1/2 resonance couplings with the decuplet

The decays listed in the particle data book, which determine the coupling constants C_{B^*} and s_c , are $N(1440) \rightarrow \Delta(1232)\pi$ and $N(1535) \rightarrow \Delta(1232)\pi$. For the sake of brevity we will only present the case of the $1/2^+$ -resonances. The parameter s_c of the $1/2^-$ octet has been determined in a completely analogous way and we just present the result. The width follows via

$$\Gamma = \frac{1}{8\pi M_{B^*}^2} |\mathbf{k}_\pi| |\mathcal{T}|^2 \quad (\text{A.1})$$

with

$$|\mathbf{k}_\pi| = \frac{1}{2M_{B^*}} \left[(M_{B^*}^2 - (M_\Delta + m_\pi)^2) (M_{B^*}^2 - (M_\Delta - m_\pi)^2) \right]^{1/2} \quad (\text{A.2})$$

being the three-momentum of the pion in the rest frame of the spin-1/2 resonance. The terms M_{B^*} and M_Δ are the masses of the octet resonance and the decuplet, respectively. For the transition matrix one obtains

$$|\mathcal{T}|^2 = \frac{C_{B^*}^2 M_{B^*}}{3F_\pi^2 M_\Delta^2} (E_\Delta + M_\Delta) \left[(M_{B^*}^2 - M_\Delta^2 - m_\pi^2)^2 - 4M_\Delta^2 m_\pi^2 \right]. \quad (\text{A.3})$$

Using the experimental value for the decay width we arrive at the central value

$$C_{B^*} = 1.35 \quad (\text{A.4})$$

where we have chosen the sign of C_{B^*} to be positive since this leads to a better fit for the Ω^- decays. From a similar calculation in the case of the $1/2^-$ -resonances one obtains

$$s_c = -0.85. \quad (\text{A.5})$$

We do not present the uncertainties in these parameters here, since for the purpose of our considerations a rough estimate of these constants is sufficient.

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Figure captions

Fig.1 Shown are the diagrams without spin-1/2 resonances in weak nonleptonic Ω^- decay. Solid and dashed lines denote ground state baryons and pseudoscalar mesons, respectively. The double line is a decuplet state. The solid square represents a weak vertex and the solid circle denotes a strong vertex.

Fig.2 Shown is the diagram with spin-1/2 resonances in the weak $\Omega^- \rightarrow \Lambda K^-$ decay. Solid and dashed lines denote the Λ and the kaon, respectively. The double line is the Ω^- and the thick line denotes the spin-1/2 resonance. The solid square represents a weak vertex and the solid circle denotes a strong vertex.

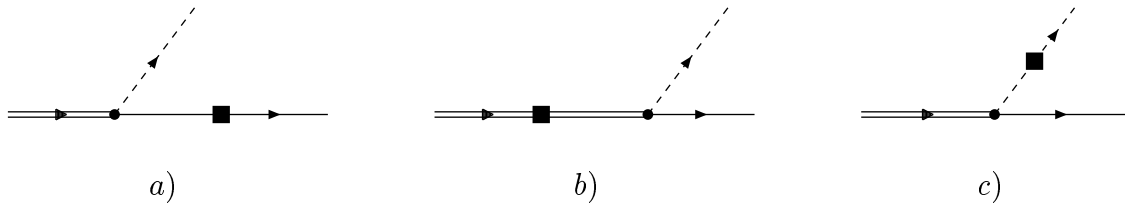


Figure 1

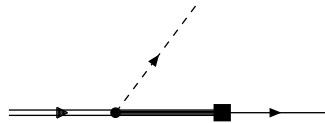


Figure 2